

The $J = \frac{3}{2}$ Θ^* partner to the $\Theta(1540)$ baryon

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Abstract

If the exotic baryon $\Theta(1540)$ is $udud\bar{s}$ with $J^P = \frac{1}{2}^+$, we predict that there is a $\overline{10}$ with $J^P = \frac{3}{2}^+$ containing a $\Theta^*(1540 - 1680)$. The width $\Gamma(\Theta^* \rightarrow KN)$ is at least a factor of three larger than $\Gamma(\Theta)$. The possibilities of $\Theta^* \rightarrow KN\pi$ or $\Theta\gamma$ via $M1$ and $E2$ multipoles are discussed.

A major plank in establishing the constituent quark model was the absence of baryons with strangeness +1. The announcement of such a particle, $\Theta(1540)$, and with a narrow width is therefore startling [1], though there is still some debate as to its existence [2, 3]. It is thus important to seek further evidence of such hadrons in order to isolate the underlying dynamics of strong QCD. We show here that if Θ is $udud\bar{s}$ with $J^P = 1/2^+$, then the correlations among QCD forces necessarily imply there be Θ^* , $J^P = 3/2^+$, which is probably only a few tens of MeV more massive.

When the proton is viewed at high resolution, as in inelastic electron scattering, its wavefunction is seen to contain configurations where its three “valence” quarks are accompanied by further $q\bar{q}$ in its “sea”. The three quark configuration is thus merely the simplest required to produce its overall positive charge and zero strangeness. The question thus arises whether there are baryons for which the minimal configuration cannot be satisfied by three quarks. The Θ would be an example; the positive strangeness requires an \bar{s} and $qqqq$ are required for the net baryon number, making what is known as a “pentaquark” as the minimal “valence” configuration.

Hitherto unambiguous evidence for such states in the data has been lacking; their absence having been explained by the ease with which they would fall apart into a conventional baryon and a meson with widths of many hundreds of MeV. It is perhaps this feature that creates the most tantalising challenge

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from the perspective of QCD: why does Θ have width below 10MeV, perhaps no more than 1MeV [4].

There is a considerable literature that recognises that ud in colour $\bar{\mathbf{3}}$ with net spin 0 feel a strong attraction [5, 6]. Such attraction between quarks in the color $\bar{\mathbf{3}}$ channel halves their effective charge, reduces the associated field energy and can be a basis of color superconductivity in dense quark matter [7]. It has been suggested [8, 9] that such correlations might even cause the S -wave combination to cluster as $[udu][d\bar{s}]$ which is the S -wave KN system, while the P -wave positive parity exhibits a metastability such as seen for the Θ and enables some contact with the Skyrme model [10] where Θ is a member of a $\bar{\mathbf{10}}$ with $J^P = 1/2^+$. There is the implied assumption that such a “diquark” may be compact, an effective boson “constituent”, which is hard to break-up. We denote this $[(ud)_0]$, the subscript denoting its spin, and the $[]$ denoting the compact quasiparticle. Ref. [9] considers the following subcluster for the pentaquark: $[(ud)_0][(ud)_0]\bar{s}$. Ref. [8] by contrast assumes that the $[(ud)_0]$ seed is attracted to a strongly-bound “triquark” $[(ud)_1\bar{s}]$.

More generally, in any pentaquark model with positive parity, angular momentum is required to annul the intrinsic negative parity of the \bar{q} . Coupling $L = 1$ with $\vec{S}_{\bar{q}}$ thus implies that both $1/2^+$ (Θ) and $3/2^+$ (Θ^*) exist. The mass gap $\Delta m(\Theta^* - \Theta)$ is determined by the strength of the spin-orbit forces within the pentaquark. No estimate of these exists in the literature to the best of our knowledge. This is the issue that we address here. We shall argue that $\Delta m(\Theta^* - \Theta) \leq m_\pi$ such that $\Theta^* \rightarrow \Theta\gamma$ transitions via $M1$ and $E2$ radiation, $\Theta^* \rightarrow KN$ and $\Theta^* \rightarrow KN\pi$ are the only allowed decay channels.

Spin orbit forces, Thomas precession and Wigner rotation effects transforming as $\langle \mathbf{L.S} \rangle$ are calculated to be individually sizable in $q\bar{q}$ and qqq states. For $q\bar{q}$ there is a cancellation among these effects arising from the short-range vector interaction (single gluon exchange) and long range confinement (assumed to be scalar) [6, 11]. For qqq baryons the situation is more subtle [11, 12] and appears to violate Galilean invariance. Ref [11, 12] showed this violation to be illusory and arises because the $\mathbf{p}=0$ frame of the two quarks experiencing the $\mathbf{L.S}$ interaction is not in general the same as the overall $\mathbf{P}=0$ frame of the qqq three-body system. Calculations performed in the $\mathbf{P}=0$ frame give the correct answer; in other frames further Wigner rotations would be required, which transform like $\mathbf{L.S}$ and contribute to the whole.

For pentaquark systems with tightly clustered *scalar* diquarks (as in refs. [8, 9]) the pattern of $\mathbf{L.S}$ and Thomas precession will differ radically from the systems $q\bar{q}$ and qqq where all pairwise interactions are between fermions. We summarise the results here and give more details elsewhere.

Spin-Orbit Splitting

We consider the spin-orbit splitting by analogy to that used with some success in the conventional meson and baryon sector. Conventionally we would consider the non-relativistic reduction of the exchange of a particle having some arbitrary propagator between two spin-1/2 quarks leading to a Breit-Fermi Hamiltonian. Such a Hamiltonian will contain the binding potentials and relativistic corrections which include spin-orbit terms. The particular form of the spin-orbit terms is fixed by the exchange propagator and the Lorentz transformation property of the vertices. Phenomenological success has come from using a combination of a scalar confining potential $V_S(r) = br$ and a vector (γ^μ) one gluon exchange with $V_V(r) = -\frac{2}{3}\frac{\alpha_s}{r}$ (between two $\mathbf{3}$'s coupled to a $\mathbf{\bar{3}}$, there is an extra factor of 2 between a $\mathbf{3}$ and a $\mathbf{\bar{3}}$ coupled to a singlet).

For quarks of mass m , the pairwise spin-orbit interaction in a vector potential takes the form [11]

$$H_{SO}^V(ij) = \frac{1}{4m^2 r_{ij}} \frac{dV_V}{dr_{ij}} (3\vec{L}_{ij} \cdot \vec{S}_{ij}^+ - \vec{K}_{ij} \cdot \vec{S}_{ij}^-) \quad (1)$$

where $\vec{L}_{ij} \equiv \vec{r}_{ij} \times (\vec{p}_i - \vec{p}_j)$; $\vec{K}_{ij} \equiv \vec{r}_{ij} \times (\vec{p}_i + \vec{p}_j)$ and $\vec{S}_{ij}^\pm \equiv (\vec{\sigma}_i \pm \vec{\sigma}_j)/2$. The analogue for a scalar potential is

$$H_{SO}^S(ij) = -\frac{1}{4m^2 r_{ij}} \frac{dV_S}{dr_{ij}} (\vec{L}_{ij} \cdot \vec{S}_{ij}^+ + \vec{K}_{ij} \cdot \vec{S}_{ij}^-) \quad (2)$$

For $q\bar{q}$ with $\mathbf{P} = 0$ one has $\mathbf{K} \equiv 0$ and there is a cancellation between V_V and V_S contributions from the $\mathbf{L.S}$ terms leading to a small net spin-orbit splitting, in line with data. For qqq the $\mathbf{K.S}$ terms appear to violate Galilean-invariance. This has caused them to be ignored in some treatments of the spin-orbit splittings of baryons [13]. The resolution of this involves the great care necessary when separating c.m. and internal coordinates for particles with spin [11] the details of which go beyond the present paper.

The essential result of refs. [11] is that the interaction in eqs.(1,2) is correct if applied in the overall rest-frame $\mathbf{P} = 0$ for an N -body system. This generalises to arbitrary systems of spinors and scalars: construct the corresponding interaction between spin-0 and spin-1/2 objects and apply the Hamiltonian in the overall $\mathbf{P} = 0$ frame.

This is our point of departure for the computation of the spin-orbit energy shifts of a pentaquark in the models of ref [8,9]. We assume that the interaction does not resolve any quark substructure within the diquark or triquark so that they can be considered to be point-like objects. Labeling the scalar by the subscript 0 and the fermion by f , the spin-orbit term so obtained has the

following form,

$$H_{SO} = \frac{\vec{\sigma} \cdot \vec{r} \times \vec{p}}{4m^2} \left(\frac{1}{r} \frac{dV_V}{dr} - \frac{1}{r} \frac{dV_S}{dr} \right) - \frac{\vec{\sigma} \cdot \vec{r} \times \vec{k}}{2mm_0} \left(\frac{1}{r} \frac{dV_V}{dr} \right), \quad (3)$$

where $\vec{r} = \vec{r}_f - \vec{r}_0$, $\vec{p} = \vec{p}_f$, $\vec{k} = \vec{p}_0$ and $\vec{\sigma}$ is the vector of Pauli spin matrices which act on the spin-1/2 state-vector.

Karliner-Lipkin model

The KL model [8] has the Θ as an effective 2-body system consisting of a scalar diquark and a spin-1/2 triquark in a relative P -wave. In their paper the masses are quoted as 720 MeV and 1260 MeV but this does not include the spin-spin interaction energy that binds them. The Θ and the D_s composite systems then have roughly the same reduced mass and are then assumed [8] to be bound by the same (QCD) dynamics. We exploit this analogy and set the diquark-triquark binding potential to be $V_V(r) = -\frac{4}{3}\frac{\alpha_S}{r}$, $V_S(r) = br$ with $b = 0.18\text{GeV}^2$ and $\alpha_S \sim 0.5$ (which gives a good fit to the hyperfine shifts of $q\bar{q}$ and qqq) [6]. In this potential the P -wave D_s mesons can be described by a variational harmonic oscillator wavefunction, $R(r) \sim r \exp -\beta^2 r^2/2$ with $\beta \sim 0.4\text{GeV}$, which reproduces the results of Godfrey and Isgur [14] to $\sim 10\%$.

With $\vec{r} = \vec{r}_{\text{tri}} - \vec{r}_{\text{di}}$ the internal momenta are $\vec{p} = \vec{p}_r$, $\vec{k} = -\vec{p}_r$ and the orbital angular momentum, $\vec{L} = \vec{r} \times \vec{p}_r$. The spin-orbit splitting term in eq.(3) becomes

$$H_{SO} = \frac{\vec{S} \cdot \vec{L}}{2m_{\text{tri}}^2} \left(\frac{4\alpha_S}{3} \left\langle \frac{1}{r^3} \right\rangle \left[1 + 2\frac{m_{\text{tri}}}{m_{\text{di}}} \right] - b \left\langle \frac{1}{r} \right\rangle \right)$$

and using $\langle L = 1 | r^{-3(-1)} | L = 1 \rangle = \frac{4}{3\sqrt{\pi}}\beta^{3(1)}$ this gives a splitting of

$$\Delta E(\Theta^* - \Theta) = \frac{1}{\sqrt{\pi}m_{\text{tri}}^2} \left(\frac{4\alpha_S}{3} \beta^3 \left[1 + 2\frac{m_{\text{tri}}}{m_{\text{di}}} \right] - b\beta \right). \quad (4)$$

Using $m_{\text{di}} = 720\text{MeV}$ and $m_{\text{tri}} = 1260\text{MeV}$ [8]

$$\Delta E(3/2 - 1/2) = (63\text{MeV})_V + (-25\text{MeV})_S = 38\text{MeV},$$

where a cancellation between vector and scalar terms is observed much as in the conventional $q\bar{q}$ and qqq states. Each term is individually small due to the large m_{di} and m_{tri} .

We investigate how much this splitting could be increased if we minimise these “constituent” masses by (i) including the subtracted hyperfine energy internal to the diquark and triquark (ii) maximising the orbital excitation energy to

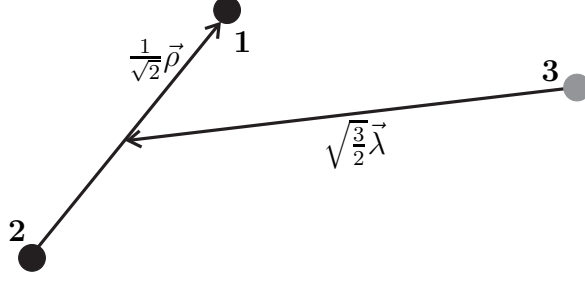


Fig. 1. 3-body variables

500MeV. If we insist that in general the sum of the diquark and triquark masses is 1540 MeV less a P -wave energy under 500 MeV, we cannot get a spin-orbit splitting greater than about 150 MeV which is still below threshold for the strong channels $\Theta^* \rightarrow \Theta\pi\pi$ or $K\Delta$ (which is forbidden for $\Theta^* \supset \overline{10}$).

Jaffe-Wilczek model

In the JW model [9] the Θ is effectively a 3-body system of two identical scalar diquarks in a P -wave and an antiquark in a relative S -wave. The color structure is anti-baryon-like $\overline{\mathbf{3}} \otimes \overline{\mathbf{3}} \otimes \overline{\mathbf{3}} = \mathbf{1}$. The Breit-Fermi Hamiltonian for this system is obtained by summing over exchanges between the three bodies pairwise. Label the diquarks 1, 2 (mass m_0); the antiquark 3 (mass m); and define (see Fig.1)

$$\begin{aligned}\vec{r}_{1,2} &= \vec{R} + \sqrt{\frac{3}{2}} \frac{m}{2m_0 + m} \vec{\lambda} \pm \frac{1}{\sqrt{2}} \vec{\rho} \\ \vec{r}_3 &= \vec{R} - \sqrt{\frac{3}{2}} \frac{2m_0}{2m_0 + m} \vec{\lambda}.\end{aligned}$$

In the Θ rest frame the internal momenta are

$$\begin{aligned}\vec{p}_{1,2} &= \frac{1}{\sqrt{6}} \vec{p}_\lambda \pm \frac{1}{\sqrt{2}} \vec{p}_\rho \\ \vec{p}_3 &= -\sqrt{\frac{2}{3}} \vec{p}_\lambda.\end{aligned}$$

Hence the one unit of orbital angular momentum between the diquarks is with respect to the $\vec{\rho}$ variable, $L_\rho = 1$ whereas the $\vec{\lambda}$ variable is in an S -wave, $L_\lambda = 0$. (This is opposite to the $L = 1$ qqq where the symmetry exposes the excitation of the λ oscillator). The spin-orbit term is then (note, exchanges between the scalar diquarks do not contribute here)

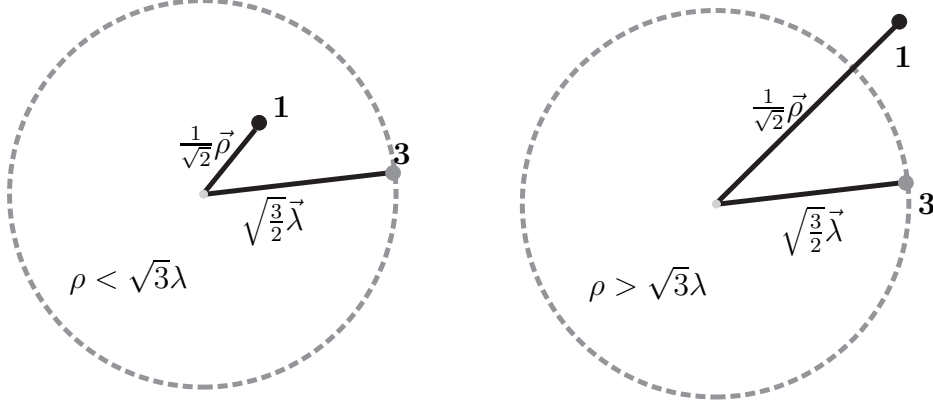


Fig. 2. after averaging over directions of $\vec{\lambda}$, the diquark feels no force when inside the sphere and feels a force in the direction of $\vec{\rho}$ when outside.

$$H_{SO} = \left\{ \left(1 + \frac{m}{m_0} \right) \Sigma_V - \Sigma_S \right\} \frac{\vec{\sigma} \cdot \vec{\lambda} \times \vec{p}_\lambda}{4m^2} + \Sigma_V \frac{\vec{\sigma} \cdot \vec{\rho} \times \vec{p}_\rho}{4mm_0} \\ + \frac{1}{\sqrt{3}} \left\{ \left(1 + \frac{m}{m_0} \right) \Delta_V - \Delta_S \right\} \frac{\vec{\sigma} \cdot \vec{\rho} \times \vec{p}_\lambda}{4m^2} + \sqrt{3} \Delta_V \frac{\vec{\sigma} \cdot \vec{\lambda} \times \vec{p}_\rho}{4mm_0}, \quad (5)$$

where $\Delta_{V,S} = \frac{1}{r_+} \frac{dV_{V,S}}{dr_+} - \frac{1}{r_-} \frac{dV_{V,S}}{dr_-}$, $\Sigma_{V,S} = \frac{1}{r_+} \frac{dV_{V,S}}{dr_+} + \frac{1}{r_-} \frac{dV_{V,S}}{dr_-}$, $\vec{r}_\pm = \vec{r}_3 - \vec{r}_{1,2}$ and $r_\pm = |\vec{r}_\pm| = \frac{1}{\sqrt{2}} \sqrt{3\lambda^2 \pm 2\sqrt{3}\vec{\lambda} \cdot \vec{\rho} + \rho^2}$. We again use vector one-gluon-exchange and now a general scalar potential to describe the binding.

The scalar interaction does not contribute to the spin-orbit splitting in this model. This is because terms featuring \vec{p}_λ are trivially zero - \vec{p}_λ acting on the $L_\lambda = 0$ wavefunction is proportional to $\vec{\lambda}$, so that the first term is $\vec{\lambda} \times \vec{\lambda}$. The $\vec{\rho} \times \vec{p}_\lambda$ term $\sim \vec{\rho} \times \vec{\lambda}$, integrating over the direction of $\vec{\lambda}$ gives something in the direction of $\vec{\rho}$ and hence $\vec{\rho} \times \vec{\rho} = 0$.

The two terms proportional to \vec{p}_ρ , which are driven by the vector interaction, are non-zero only when $\sqrt{3}\lambda \leq \rho$ which is a consequence of Gauss's law. They are proportional to the force on the diquarks due to the antiquark which is only non-zero when the diquark is outside the spherical shell obtained by averaging over directions of $\vec{\lambda}$ with fixed λ - see Fig.2. This is rather similar to the baryon model considered in [15]. Thus although there is no cancellation between vector and scalar in this model, the spatial restriction to the spherical shell defined by ρ enfeebles the total $\mathbf{L} \cdot \mathbf{S}$ contribution here.

Specifically, after integration over the direction of $\vec{\lambda}$ these terms combine to give a spin-orbit term for $\rho > \sqrt{3}\lambda$,

$$\frac{2\sqrt{2}}{mm_0} \frac{2\alpha_S}{3} \frac{1}{\rho^3} \vec{S} \cdot \vec{L}_\rho,$$

and a splitting

$$\Delta E(3/2 - 1/2) = \frac{2\sqrt{2}\alpha_S}{mm_0} \left\langle \frac{1}{\rho^3} \theta(\rho - \sqrt{3}\lambda) \right\rangle$$

where only the vector interaction contributes.

Jaffe and Wilczek do not explicitly state the masses of their diquarks. The standard **S.S** interaction would give $m_0 \sim 500\text{MeV}$, but this is hard to confront with a $\Theta(1540)$ containing two diquarks, together with m_s and a P -wave excitation energy (ω) as well. Further, one should ensure that $[ud]_0[ud]_0[ud]_0$ with two P -waves, is not more stable than the deuteron. This requires that $\omega \gtrsim 450\text{MeV}$. Such a result is at least consistent with the (anti)baryon spectrum, which has the same internal colour arrangement and by assumption similar binding dynamics for which the energy gap between $(m_N + m_\Delta)/2$ and the negative parity $N^*(1520 - 1750)$ is $\sim 500\text{MeV}$. This leaves $\sim 1 - 1.1\text{GeV}$ to be shared between two diquarks and the \bar{s} . With $m_s \sim 450\text{MeV}$ we thus assign a mass of $\sim 350\text{MeV}$ to each diquark. This is the minimum we can tolerate, and even with this we shall find that the $\Theta^* - \Theta$ mass gap is only tens of MeV; any larger mass for the diquarks would reduce it even more.

We can make an estimate of the $\Theta^{(*)}$ spatial wavefunction from baryon measurements. The $\Lambda - \Sigma$ splitting can be used to find the harmonic oscillator parameter, $\alpha_\rho \approx 400\text{MeV}$;

$$M(\Sigma^0) - M(\Lambda^0) \approx 77\text{MeV} = \frac{16\pi\alpha_S}{9m} \left(\frac{1}{m} - \frac{1}{m_s} \right) \frac{\alpha_\rho^3}{\pi^{3/2}}.$$

This gives a level spacing $\omega(L = 0, L = 1) = \alpha_\rho^2/m \sim 485\text{MeV}$ in good agreement with the data [16].

The radial integrals can be performed with the approximate harmonic oscillator wavefunctions giving

$$\Delta E(3/2 - 1/2) = \frac{2\sqrt{2}\alpha_S}{mm_0} \left(\frac{4}{3\sqrt{\pi}} \frac{k^3}{(3 + k^2)^{3/2}} \right) \alpha_\rho^3 \quad (6)$$

where $k = \left(\frac{3m}{2m_0 + m} \right)^{1/4}$. With $\alpha_\rho \sim 400\text{MeV}$ this gives $\Delta E \sim 35, 65\text{MeV}$ with $m_0 = 500, 300\text{MeV}$ respectively.

In all cases considered here, the spin-orbit excitation energy is plausibly small on the scale of the mass gap between scalar and vector diquark, which is expected to be $O(200)\text{MeV}$. That this is so is an explicit assumption here, and implicitly assumed in ref. [9] for the required correlations to obtain. Thus we conclude that the first Θ^* will be the **L.S** state discussed here; states where the clusters are internally excited, if any, will be at higher masses. However, if

the clusters can be resolved, there is the possibility of dilution of the present effects, or of introducing non-zero contributions from tensor forces for example. It cannot be excluded that these could conspire to reduce the net mass gap, even that the Θ^* and Θ are degenerate within the present resolution of the experimental data.

Other states

Ref [9] consider $[(ud)_0][(ud)_0]\bar{Q}$ with $Q \equiv u, d$. These are expected to lie $O(100 - 150)\text{MeV}$ below the Θ , and may be identified with the $P_{11}(1440)$ ‘‘Roper’’ resonances [16]. The **L.S** forces then imply a $3/2^+$ partner at $\sim 1.5 - 1.6\text{GeV}$, for which there is no evidence [16]. However, this $\mathbf{\bar{10}}\text{-}\mathbf{8}$ mixture can couple to conventional qqq (nucleon) by $q\bar{q}$ annihilation, which may explain both its large width and possibly lead to large mass shifts, and unobservably large widths of any $3/2^+$ partner. For the $[(ud)_0][(us)_0]\bar{s}$ state, the mass gap to the $3/2^+$ partner is predicted to be $\sim O(10)\text{MeV}$. This tantalisingly is in accord with the $P_{11}(1710) - P_{13}(1720)$ pair.

The mixed messages here may indicate that the concept of constituent pentaquarks is meaningful only for manifestly exotic combinations. The dynamics of **L.S** may therefore be probed also by the exotic Ξ^+ , Ξ^{--} and the heavy flavour analogues $[ud][ud]\bar{Q}$, with $Q \equiv c, b$. We predict $\Delta m(\Xi^* - \Xi) \sim 30 - 50\text{MeV}$. The reason that this is similar to $\Delta m(\Theta^* - \Theta)$ is due to the effect of the heavier s mass being ‘‘diluted’’ within clusters, and the \bar{s} being replaced by the lighter $\bar{d}(\bar{u})$. The splitting scales as the inverse constituent mass for large m_Q (eqns.4, 6) and hence the splittings for $\Theta_c^* - \Theta_c$ are at most a few MeV.

Decays of Θ^*

The ‘‘natural’’ width of a P -wave Θ resonance 100MeV above NK threshold is of order 200MeV [9, 17]. However, this has not yet taken into account any price for recoupling colour and flavour-spin to overlap the $(ud)(ud)\bar{s}$ onto NK colour singlets such as uud and $d\bar{s}$. In amplitude, starting with the Jaffe-Wilczek configuration, the colour recoupling costs $\frac{1}{\sqrt{3}}$ and the flavour-spin to any particular channel (e.g. K^+n) costs a further $\frac{1}{2\sqrt{2}}$. If in addition we suppose that spatially the constituents then fall apart, only the $L_z = 0$ piece of the wavefunction contributes, which implies a further suppression in amplitude of $\sqrt{\frac{1}{3}}$ for the $L = 1 \otimes S = \frac{1}{2} \rightarrow J = \frac{1}{2}$. Thus a total suppression in rate of up to two orders of magnitude may be accommodated in such pictures.

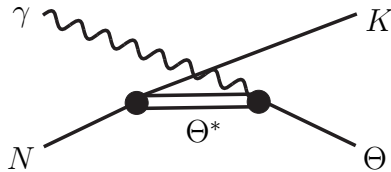


Fig. 3. u-channel production diagram.

For the Θ^* similar arguments obtain, though Γ_{TOT} is now expected to be larger as (i) the spatial Clebsch is now $\sqrt{\frac{2}{3}}$ instead of $\sqrt{\frac{1}{3}}$; (ii) the phase space $\sim q^3$ gives a factor 1.8 – 3.5. So $\Gamma(\Theta^*(1600)) \sim 3.6\Gamma(\Theta)$, rising to $\sim 7\Gamma(\Theta)$ if $m \sim 1700\text{MeV}$. However, if the decay involves tunneling, which is exponentially sensitive to the difference between the barrier height and the kinetic energy of the state, this could significantly enhance the width of the Θ^* [17].

With the mass gaps as predicted here, $\Theta^* \rightarrow \Theta\pi\pi$ is kinematically forbidden; the possibility of $\Theta^* \rightarrow KN\pi$ emerges, but the phase space for a non-resonant three body decay may prevent this being a large branching ratio (unless the $N\pi$ is enhanced by the tail of a pentaquark $P_{11}(1440)$). The $b.r(\Theta^* \rightarrow \Theta\gamma)$ is $\leq 10^{-3}$ due to the restricted phase space. However the ratio of $M1$ amplitudes for $\gamma\Theta \rightarrow \Theta^*$ and $\gamma N \rightarrow \Delta$, with momentum factors removed, is $\frac{m_u}{3} \left(\frac{1}{m_s} + \frac{1}{m_0} \right) \sim 0.4$ and so it is possible that this feeds a significant $\gamma N \rightarrow \Theta K$ via Θ^* exchange in the u -channel (Fig 3). Whereas the $\Delta \rightarrow N\gamma$ in the constituent quark model has $E2 \equiv 0$ due to the $L = 0$ internal structure of these baryons, the $\Theta^* \rightarrow \Theta\gamma$ has $E2$ arising from the internal orbital degree of freedom.

In summary pentaquark models of $J^P = 1/2^+$ $\Theta(1540)$ imply there be a copy of the $\overline{\mathbf{10}}$ (and $\mathbf{8}$) containing Θ^* with $J^P = 3/2^+$ and within tens of MeV of it. We advocate searching for this in KN or $KN\pi$ final states. There is also the possibility that Θ^* is actually the 1540 state already observed, and that the true $1/2^+$ state lies around 1500MeV, or that the two states are degenerate within the present resolution of the data. Either of these could explain the conundrum of why there is no clear sign of a more prominent narrow structure in the KN spectrum above 1540MeV. The possibility that the Θ^* is broad [17] might also explain the present data. An alternative possibility is that the observed Θ is the decay product of some directly produced particle. In the photoproduction experiments this could be via the Θ^* in the u -channel. The spectrum of the Skyrme model as described in e.g. [18] has no place for such a $J^P = 3/2^+$ $\overline{\mathbf{10}}$ multiplet. As such our prediction could be an interesting test of models [19].

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